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Axino-induced baryogenesis

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We consider the possibility that the baryon asymmetry is generated at low energies as a consequence of the axino decay. We analyse models in which the axino, the superpartner of the axion, is heavy and decays into gluinos at temperatures $T \sim 1$ GeV. If CP and B violating couplings for the quark superfields are included, the subsequent decay of these out of equilibrium gluinos can act as seeds for baryogenesis. The required amount of CP violation is well consistent with the bounds on the electric dipole moment of the neutron and the mechanism works even for low reheating temperatures after inflation ($T_{RH} \gtrsim 10^4$ GeV).

I. INTRODUCTION

The standard explanation for the origin of the observed baryonic asymmetry $\eta \equiv n_B/n_\gamma \sim 10^{-10}$ requires B and CP violating processes occurring out of equilibrium during the early stages of the evolution of the universe.[1] These so-called Sakharov conditions are naturally met in some grand unified theories, which contain heavy Higgs (and vector) bosons with the desired couplings that may be left out of equilibrium by the fast expansion of the universe, producing the baryon asymmetry when they decay.

The required conditions are in principle also attainable in the standard model, that allows for anomalous B violation, although in this case the out of equilibrium situation is difficult to realize at the low temperatures involved. Since the anomalous electroweak B violating (and $B-L$ conserving) processes are in equilibrium up to temperatures $T \sim m_W/\alpha$, even if they happen not to be relevant for the generation of η, they are important because they may erase a baryon asymmetry generated previously (e.g. in a GUT scenario), unless $B-L \neq 0$ initially.[2] Furthermore, if the anomalous processes are combined with other sources of low energy L or B violation that are present in many extensions of the standard model (Majorana masses, R parity violating couplings, etc), they can eliminate any preexisting combination of B and L.[3]

Another possible problem for the GUT baryogenesis is that if the universe undergoes a period of inflation, as is required to solve the problems of flatness, horizon, monopoles, etc., the reheating temperature after the end of inflation may not be high enough to allow the heavy bosons to be thermally produced again ($T_{RH} \lesssim 10^{12}$GeV), in which case one is forced to devise an alternative mechanism leading to baryogenesis at low energies. In addition to the electroweak baryogenesis,[2,4] some other scenarios for this have been proposed. In particular, the supersymmetric extension of the stan-

dard model has been a playing ground for these attempts. Baryon violation can be obtained in these models by including the superpotential couplings $g_{ijk}u_i^c d_j^c d_k^c$, that are consistent with gauge and supersymmetry invariance and, if not combined with other L violating Yukawa couplings, are compatible with proton stability. Many new sources of CP violation are also usually present in supersymmetric models,[5] and the source of non-equilibrium has been searched in late-decaying particles. For example, in ref. [6] out of equilibrium squarks are produced by the decay of the inflaton itself and ref. [7] exploits the gauginos that result from the late gravitino decays. These out of equilibrium particles can then generate the baryon asymmetry in their subsequent decay. In particular, the gravitino induced baryogenesis takes advantage of the possible overabundance of gravitinos, that results from their very early decoupling at the Planck epoch and their survival until the decay time $\tau_D \sim 8\pi M_P^2/m_{3/2}^3 \simeq (10\ \mathrm{TeV}/m_{3/2})^3$sec. To avoid disrupting the nucleosynthesis predictions of the light elements abundance, this scheme requires gravitino masses $m_{3/2}$ in excess of ~ 50 TeV. One possible inconvenience of this model is that if inflation dilutes the primordial gravitinos, their regeneration up to a sufficient amount, that is due to processes with gravitational strength, requires extremely large reheating temperatures, greater than $\sim 10^{15}$ GeV.

In this paper we want to consider a new possibility that appears when one combines supersymmetry with the Peccei Quinn solution to the strong CP problem. In this frame, besides the axion the theory contains its superpartner, the axino $\tilde{a}$, which can have very important cosmological implications.[8] Since the axino couplings are suppressed by the PQ symmetry breaking scale F, that is constrained to be $10^{10} \div 10^{13}$GeV,[9] axinos decouple at very early times surviving in large amounts. If axinos are stable, they would overclose the universe unless $m_{\tilde{a}} \lesssim 2$ keV[8] and hence axinos heavier than this bound should be unstable. We will consider henceforth the supersymmetrized version of the invisible axion model of Kim [10], in which axinos

are naturally heavy, with a mass that can be comparable to the gravitino mass. If $m_{\tilde{a}} \sim 1$ TeV axinos will mainly decay into a gluon gluino pair with a lifetime

$$\tau \simeq \frac{9\pi^3}{8\alpha_s^2}\frac{F^2}{m_{\tilde{a}}^3} \sim 10^{-6}\left(\frac{F}{10^{12}\mathrm{GeV}}\right)^2\left(\frac{m_{\tilde{a}}}{1\mathrm{TeV}}\right)^{-3} \mathrm{sec}. \tag{1}$$

(This is similar to the $\tilde{\gamma} \rightarrow \tilde{a}\gamma$ lifetime computed in ref. [12] but including the color factor.) For the allowed values of F these decays correspond to temperatures around 1 GeV and thus do not interfere with nucleosynthesis at all, and one is tempted to use them as the source for the out of equilibrium gluinos that can act as seeds for baryogenesis. Furthermore, since axinos are regenerated much more efficiently than gravitinos, this mechanism works also for much smaller reheating temperatures than those required in the gravitino induced baryogenesis of ref. [10]. As it will also turn out, the CP violation necessary to account for sufficient baryon asymmetry generation can be obtained for a wide range of the parameter space in a way compatible with the experimental constraints such as the induced electric dipole moment of the neutron. One should also note that in the scenario to be discussed here the natural dark matter constituent is the axion itself, which is one of the preferred Cold Dark Matter candidates.

II. AXINO COSMOLOGY

We will consider the supersymmetric version of the heavy quark invisible axion model proposed by Kim.[10] In this model, a heavy color triplet Q is introduced which can be taken to be an $SU(2) \times U(1)$ singlet. Its right-handed component Q_R carries PQ charge $+1$, and the left-handed component Q_L carries -1. The relevant Yukawa coupling is given by

$$\mathcal{L}_Y = -f\bar{Q}_L\sigma Q_R - h.c. \,, \tag{2}$$

where σ is the complex scalar field singlet of $SU(3) \times SU(2) \times U(1)$ in which phase

resides the axion field. The PQ symmetry is realized for a PQ charge -2 for σ.

When the PQ symmetry is spontaneously broken at a scale $F = \langle\sigma\rangle$, the quarks Q attain a mass $f\langle\sigma\rangle$. The main contribution to the axino mass comes from the one loop diagram represented in figure 1, that leads to[13]

$$m_{\tilde{a}} = \frac{3}{16\pi^2} f^2 A m_{3/2} \simeq 1 \text{ TeV} \left(\frac{f^2 m_{3/2}}{10 \text{ TeV}}\right) \tag{3}$$

where A is a number of order one. For $f \simeq \mathcal{O}(1)$, as we will assume, the axino is naturally heavy and to avoid overclosing the universe it should be unstable.

In this model, the axion mainly interacts with ordinary matter through a loop of heavy quarks Q coupling to two gluons. This vertex can be written as[11]

$$\mathcal{L}_{aGG} = \frac{\alpha_s}{16\pi\ F}\ a\ \epsilon_{\mu\nu\rho\delta} G^{c\mu\nu} G^{c\rho\delta}. \tag{4}$$

The corresponding supersymmetric vertex of the axino-gluon-gluino interaction is given by

$$\mathcal{L}_{\tilde{a}G\tilde{G}} = \frac{\alpha_s}{8\pi\ F}\ \tilde{a}\gamma_5\sigma^{\mu\nu}\tilde{G}^c G^c_{\mu\nu}. \tag{5}$$

If gluinos are lighter than the axino, this coupling allows for the decay $\tilde{a} \to g\tilde{g}$ with the lifetime given in eq. (1).

In order to study the cosmological evolution of axinos, one should follow their thermal history. At high temperatures the axinos are kept in thermal equilibrium through the gluon mediated reactions $q\bar{q} \leftrightarrow \tilde{a}\tilde{g}$ and $q\tilde{g} \leftrightarrow q\tilde{a}$, which proceed at a rate

$$\Gamma \sim \frac{\alpha_s^3}{16\pi F^2} n(T), \tag{6}$$

where $n(T) = 2\zeta(3)T^3/\pi^2$ is the number density of a relativistic specie at temperature T. When Γ becomes smaller than the expansion rate of the universe H, the interactions are no more effective to maintain axinos in chemical equilibrium and they

decouple. This happens at a temperature

$$T_D \simeq 16\pi^2 \left(\frac{8\pi g_*}{90}\right)^{1/2} \frac{F^2}{M_P \alpha_s^3} \sim 10^{11}\mathrm{GeV} \left(\frac{F}{10^{12}\mathrm{GeV}}\right)^2, \tag{7}$$

where g_* is the effective number of degrees of freedom, that for temperatures higher than the superpartner masses is $g_* = 930/4$.

At temperatures larger than T_D, the axino number density equals the photon one. After the decoupling both are equally diluted by the universe expansion; however successive annihilation of species increase the number density of photons leaving unchanged that of axinos. Thus, at a given temperature $T < T_D$, and before the axinos decay, the ratio of the number density of axinos to photons is given by

$$\frac{n_{\tilde{a}}}{n_\gamma}(T) = \frac{g_*(T)}{g_*(T_D)}. \tag{8}$$

The picture may be changed if the universe underwent a period of inflation. The exponential expansion dilutes the initial number density of particles to negligible amounts. When the inflaton decays, reheating the universe, a new thermal bath of particles is created. If the reheating temperature T_{RH} is larger than T_D, the thermal bath will contain also axinos and the previous scheme holds. However, if the universe reheats to a temperature smaller than T_D, the density of axinos will never reach the thermal equilibrium one. In this case, we can compute the amount of regenerated axinos integrating the equation

$$\frac{dn_{\tilde{a}}}{dt} + 3\frac{\dot{R}}{R} n_{\tilde{a}} = \Gamma n(T), \tag{9}$$

where R denotes the scale factor of the universe. The resulting ratio of the number density of axinos to photons for $T < T_{RH}$ is

$$\frac{n_{\tilde{a}}}{n_\gamma}(T) = \frac{g_*(T)}{g_*(T_{RH})} \left(\frac{g_*(T_D)}{g_*(T_{RH})}\right)^{1/2} \frac{T_{RH}}{T_D}, \tag{10}$$

where we have introduced T_D from eq. (7).

When axinos become non-relativistic ($T \lesssim m_{\tilde{a}}$), the energy density $\rho_{\tilde{a}} = m_{\tilde{a}} n_{\tilde{a}}$ associated to the axinos decreases with temperature as T^3, while that of relativistic particles as T^4. Thus, the axinos contribution to the total energy increases as the universe expands and they may become dominant before they decay.

The temperature of the universe just after the decay process occurs, T_r, can be computed equating τ to H^{-1}. As the axinos decay into relativistic particles, assuming that the reheating is instantaneous, we obtain using eq. (1)

$$T_r = \left(\frac{90}{8\pi g_*(T_r)}\right)^{1/4} \frac{2\sqrt{2}\alpha_s}{3\pi^2} \frac{m_{\tilde{a}}^{3/2} M_P^{1/2}}{F}. \tag{11}$$

(actually the universe does not reheat but just cools more slowly during the axino decay, but this treatment is satisfactory for our purposes.)

If the baryogenesis results from the decay of the axinos, with each axino decay giving rise to ϵ baryons, then the ratio of baryons to photons at present ($T_0 \simeq 2.7^\circ K$) is given by

$$\eta \simeq \epsilon \frac{g_*(T_0)}{g_*(T_r)} \frac{n_{\tilde{a}}(\tau)}{n_\gamma(T_r)}, \tag{12}$$

where $g_*(T_0)/g_*(T_r)$ takes into account the increment of the number of photons with respect to that of baryons due to the annihilation of species after the axino decay. The baryon asymmetry can be written using eqs. (12), (1) and (11) as

$$\eta \simeq \epsilon \frac{g_*(T_0)}{g_*(T_r)} \frac{\sqrt{2}\pi\alpha_s}{8\zeta(3)} \left(\frac{8\pi g_*(T_r)}{90}\right)^{3/4} \frac{\sqrt{M_P\, m_{\tilde{a}}}}{F} \frac{\rho_{\tilde{a}}(\tau)}{\rho_{tot}(T_r)}. \tag{13}$$

The last factor measures the amount that the axinos contribute to the total energy when they decay. If they are dominant, it is of order one. This turns out to be the case when there is no inflation (or if the reheating temperature is larger than the decoupling one) and for values of the PQ scale F larger than $\sim 10^{11}$ GeV. For smaller values, the axino contributes

$$\frac{\rho_{\tilde{a}}}{\rho_{tot}} \simeq 0.5 \left(\frac{F}{10^{11}\ \mathrm{GeV}}\right) \left(\frac{m_{\tilde{a}}}{1\ \mathrm{TeV}}\right)^{-1/2}. \tag{14}$$

In the case that the reheating temperature after inflation is smaller than T_D, the right hand side of eq. (14) should be multiplied by $[g_*(T_D)/g_*(T_{RH})]^{3/2}\ T_{RH}/T_D$.

III. BARYON ASYMMETRY

In order to generate a non vanishing ϵ, we will consider the same mechanism that was proposed in ref. [7] to produce the baryon asymmetry out of gravitino decays. For this we start from a general softly broken supergravity lagrangian requiring lepton number conservation. This allows for the B violating superpotential couplings of the $SU(2)$ quark singlet superfields

$$g_{ijk}\ u_i^c d_j^c d_k^c, \tag{15}$$

with i, j, k being generation indices and for definiteness we will take only $\alpha_{332} \equiv g_{332}^2/4\pi$ to be non vanishing (stringent constraints apply to the couplings involving the first generation because of their effects on heavy nuclei decay and $n\ \bar{n}$ oscillations[6]). The relevance of this term is twofold. First, it allows for CP violation in the gluino decays, because the interference between the second diagram of fig. 2, that involves the trilinear scalar interaction $g_{332}Am_{3/2}\tilde{t}^c\tilde{b}^c\tilde{s}^c$, and the tree level term is a source of CP violation if $\mathrm{Im}(Am_{3/2}m_{\tilde{g}}) \neq 0$. The two diagrams of fig. 2 will then generate more $t\ \bar{\tilde{t}}$ than $\bar{t}\ \tilde{t}$ pairs (or viceversa, depending on the CP violating phase) in the gluino decays. As was estimated in ref. [7], the resulting asymmetry $\Delta\Gamma_{\tilde{g}} \equiv \Gamma(\tilde{g} \to t\bar{\tilde{t}}) - \Gamma(\tilde{g} \to \bar{t}\tilde{t})$ is

$$\frac{\Delta\Gamma_{\tilde{g}}}{\Gamma_{\tilde{g}}} \simeq \frac{\alpha_{332}}{4}\ \frac{Im(Am_{3/2}m_{\tilde{g}})}{|m_{\tilde{g}}|^2} G(\frac{m_{\tilde{q}}^2}{m_{\tilde{g}}^2}), \tag{16}$$

with

$$G(x) = \theta(1-x)\left(1 - \frac{x}{(1-x)^2}\ln\left(x + \frac{1}{x} - 1\right)\right). \tag{17}$$

The other important aspect of the couplings in eq. (15) is that they allow for the $\Delta B = \pm 1$ decay of the (anti)squarks produced in the gluino decay. Hence, each

gluino will give rise either to 3 quarks or 3 antiquarks, and the imbalance between these two channels, resulting from the CP violation, leads to a baryon asymmetry per axino decay $\epsilon \simeq \Delta\Gamma_{\tilde{g}}/\Gamma_{\tilde{g}}$.

In order that the generated baryon asymmetry is not erased by squark mediated processes such as $qq \to \bar{q}\chi$, we will take the neutralinos χ to be heavy ($m_\chi > 100$ GeV), so that these B violating processes are very suppressed after the axino decay. Note that in this model the neutralinos decay through $\chi \to 3q$ or $\chi \to 3\bar{q}$, and the lightest supersymmetric particle is not stable.

Looking back to eq. (13), we see that if the axinos contribute significantly to the energy density at the time they decay (this holds for $T_{RH} \gtrsim T_D$), the value of ϵ that is required to produce $\eta \sim 10^{-10}$ is only $\epsilon \sim 10^{-8} \div 10^{-7}$. One can check that the required amount of CP violation does not induce at the same time an electric dipole moment for the neutron d_n in conflict with observations. In fact, taking for instance $\alpha_{332} = 0.01$, $m_{\tilde{g}} = 300$ GeV and $m_{\tilde{q}} = 100$ GeV, the electric dipole moment generated by the phase in $Am_{\tilde{g}}$ turns out to be, using the results of ref. [14], four orders of magnitude below the experimental upper bound $|d_n| \lesssim 10^{-25}e$ cm.[15] This also allows to achieve the required values of η in scenarios with lower reheating temperatures, just by invoking larger values of ϵ that can be easily obtained. For example, in the 'extreme' case $\alpha_{332} \simeq 0.1$ and with and induced $d_n \simeq 10^{-25}e$ cm, T_{RH} may be as low as 10^4 GeV for $F \simeq 10^{11}$ GeV.

As a summary, we have considered supersymmetric invisible axion models in which axinos are heavy, as can happen in the Kim model. Assuming that the supersymmetric masses satisfy $m_{\tilde{q}} < m_{\tilde{g}} < m_{\tilde{a}} \simeq 1$ TeV we have shown that axinos decay at a temperature ~ 1 GeV, producing out of equilibrium gluinos. Introducing CP and B violating couplings these decays can act as seeds for baryogenesis. The observed

baryon asymmetry can be generated even with small amounts of CP violation and also for low reheating temperatures after inflation.

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Figure captions:

Fig. 1: One-loop contribution to the axino mass.

Fig. 2: Tree and one-loop diagrams contributing to the gluino decay. Their interference is the source of CP violation leading to $\Gamma(\tilde{g} \to t\bar{\tilde{t}}) \neq \Gamma(\tilde{g} \to \bar{t}\tilde{t})$.

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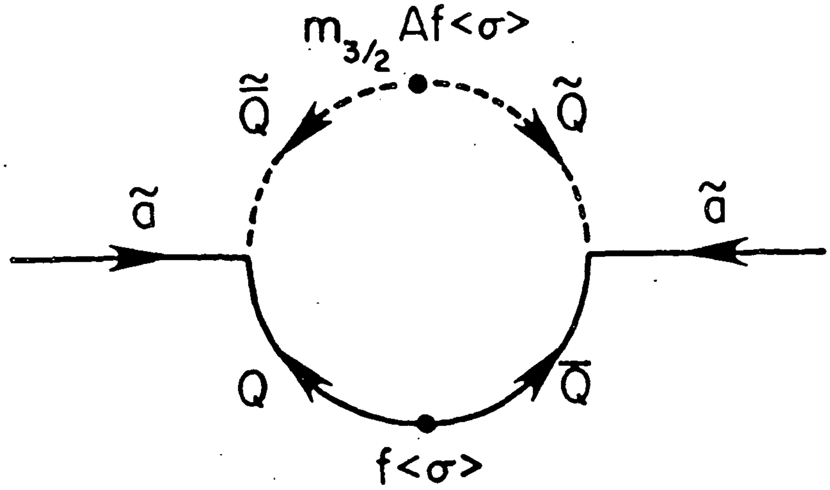
$m_{3/2} A f\langle\sigma\rangle$
$\tilde{\bar{Q}}$
$\tilde{Q}$
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FIG. 1

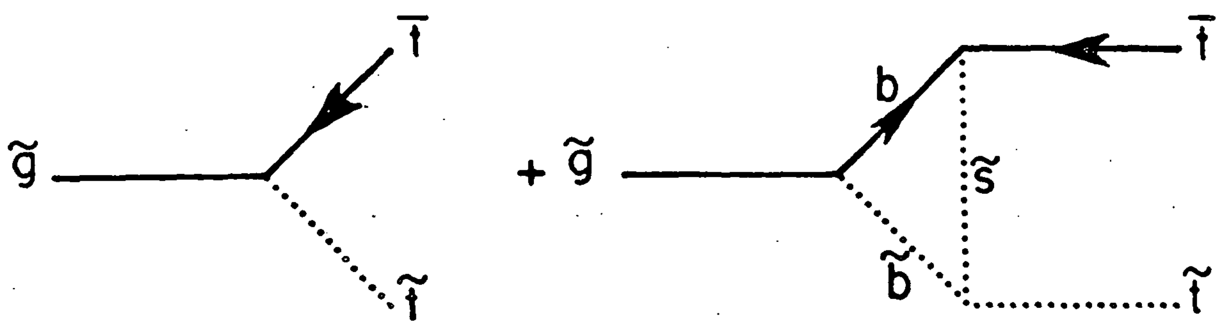
$\tilde{g}$
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FIG. 2